

# VDM: A model for Vector Dark Matter

Yasaman Farzan<sup>\*</sup> and Amin Rezaei Akbarieh<sup>†</sup>

<sup>\*,†</sup>School of physics, Institute for Research in Fundamental Sciences (IPM)

P.O.Box 19395-5531, Tehran, Iran

<sup>†</sup>Department of Physics, Sharif University of Technology

P.O.Box 11155-9161, Tehran, Iran

## Abstract

We construct a model based on a new  $U(1)_X$  gauge symmetry and a discrete  $Z_2$  symmetry under which the new gauge boson is odd. The model contains new complex scalars which carry  $U(1)_X$  charge but are singlets of the Standard Model. The  $U(1)_X$  symmetry is spontaneously broken but the  $Z_2$  symmetry is maintained, making the new gauge boson a dark matter candidate. In the minimal version there is only one complex scalar field but by extending the number of scalars to two, the model will enjoy rich phenomenology which comes in various phases. In one phase, CP is spontaneously broken. In the other phase, an accidental  $Z_2$  symmetry appears which makes one of the scalars stable and therefore a dark matter candidate along with the vector boson. We discuss the discovery potential of the model by colliders as well as the direct dark matter searches.

## 1 Introduction

Although in the recent decades overwhelming evidence for the presence of dark matter has been accumulated by astrophysical and cosmological observations, discovering the nature of Dark Matter (DM) is still one of the open questions before particle physicists. For example, we still do not know what

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<sup>\*</sup>E-mail: yasaman@theory.ipm.ac.ir

<sup>†</sup>E-mail: am-rezaei@physics.sharif.ir

is the spin of the DM candidate. Complex and real scalars as well as Dirac and Majorana fermions as the DM candidates have been extensively studied in the literature. Vector boson as the dark matter candidate has only recently received attention [1, 2, 3, 4, 5]. Although the vector boson playing the role of DM does not necessarily need to be a gauge vector boson, in most of the scenarios employing a vector boson as dark matter, it is taken to be the gauge boson of a new gauge symmetry. In [6] and [7], a new non-Abelian gauge symmetry is introduced in such a way that one [6] or all [7] of the gauge bosons play the role of the DM. Abelian gauge boson dark matter has been studied in the context of extra large dimension [8], the little Higgs model [9] and the linear sigma model [10].

In this paper, we introduce a simple model within which the Abelian gauge boson plays the role of the dark matter; *i.e.* the gauge group of the standard model is extended to  $SU(3) \times SU(2) \times U(1)_Y \times U(1)_X$  such that all the SM particles are neutral under the  $U(1)_X$ . The model also contains new complex scalar fields which are singlet under the color and electroweak symmetry but transform under the new  $U(1)_X$  symmetry. We will first introduce a model with a single complex scalar and then extend it to two complex scalar bosons. The extension results in a rich phenomenology. The model also has two  $Z_2$  symmetries under which the new vector boson is odd. In other words, the symmetry of the model becomes  $U(1)_X \times Z_2 \times Z_2$  times the SM symmetries. One of the new scalar fields receives a vacuum expectation value which breaks the new  $U(1)_X \times Z_2 \times Z_2$  symmetry into a remnant  $Z_2$  symmetry that protects the DM candidate against decay. After the gauge symmetry breaking, one of these scalars mixes with the standard model Higgs. This mixing is the only portal between the dark sector and the standard model particles. From the point of view of direct detection, the standard results for Higgs portal scenario [7] applies to this model, too. However, as we shall see, depending on the parameter range, new annihilation modes can also be open which do not involve the SM Higgs boson.

The paper is organized as follows: In section 2, we introduce the model and shortly discuss its various phases. In section 3, we discuss the annihilation of dark matter pair. In section 4, we discuss the dark matter direct and collider searches. In section 5, we discuss the lower bounds on the coupling with the Standard Model (SM) Higgs. The results are summarized in section 6.

## 2 The model

This model is based on a new Abelian gauge symmetry, under which the standard model particles are all neutral. The gauge boson of the new  $U(1)_X$  symmetry is denoted by  $V_\mu$ . We impose a  $Z_2$  symmetry under which the SM particles are all even but  $V_\mu$  is odd. As a result, the kinetic mixing between  $V_\mu$  and the hyper-charge gauge boson is forbidden by this  $Z_2$ . In section 2.1, we introduce the minimal version with a single complex scalar. The possibility is briefly discussed in [1], too. In section 2.2, we extend the model to include two complex scalars.

### 2.1 Minimal model with a single complex scalar

In this model, we include a singlet scalar  $\Phi = (\phi_r + i\phi_i)/\sqrt{2}$  which is charged under  $U(1)_X$ . The Lagrangian of the scalars and the new vector boson is

$$\mathcal{L} = D_\mu \Phi D^\mu \Phi - \frac{V_{\mu\nu} V^{\mu\nu}}{4} - V(\Phi, H), \quad (1)$$

where  $V_{\mu\nu} = \partial_\mu V_\nu - \partial_\nu V_\mu$ ,  $D_\mu = \partial_\mu - ig_V V_\mu$  and

$$V = -\mu_\phi^2 |\Phi|^2 - \mu^2 |H|^2 + \lambda_\phi |\Phi|^4 + \lambda |H|^4 + \lambda_{H\phi} |\Phi|^2 |H|^2. \quad (2)$$

Notice that the  $U(1)_X$  symmetry implies that the Lagrangian is invariant under the  $Z_2^{(A)} \times Z_2^{(B)}$  symmetry defined as follows:

$$Z_2^{(A)} : V_\mu \rightarrow -V_\mu, \quad \Phi \rightarrow \Phi^*$$

and

$$Z_2^{(B)} : V_\mu \rightarrow -V_\mu, \quad \Phi \rightarrow -\Phi^*,$$

where the rest of the fields are even.

$H$  and  $\Phi$  receive VEVs breaking respectively the electroweak and  $U(1)_X$  symmetries. Going to the “unitary” gauge, the imaginary component of  $\Phi$  can be absorbed as the longitudinal component of  $V_\mu$ . In this gauge, we can write

$$\Phi = \frac{\phi_r + v_r}{\sqrt{2}} \text{ and } H = \begin{pmatrix} 0 \\ \frac{h+v}{\sqrt{2}} \end{pmatrix} \quad (3)$$

where

$$\begin{aligned} v^2 &= \frac{4\lambda_\phi\mu^2 - 2\lambda_{H\phi}\mu_\phi^2}{4\lambda\lambda_\phi - \lambda_{H\phi}^2} \\ v_r^2 &= \frac{4\lambda\mu_\phi^2 - 2\lambda_{H\phi}\mu^2}{4\lambda\lambda_\phi - \lambda_{H\phi}^2}. \end{aligned} \quad (4)$$

The conditions for successful spontaneous symmetry breaking are  $v^2 > 0$  and  $v_r^2 > 0$ . Notice that while  $Z_2^{(B)}$  is broken, the  $Z_2^{(A)}$  symmetry still persists making  $V_\mu$  a stable particle and therefore a dark matter candidate. The mass of  $V_\mu$  is given by

$$m_V = g_V v_r.$$

$\phi_r$  mixes with  $h$  with the following mixing matrix

$$\frac{1}{2}[\phi_r \ h] \begin{bmatrix} 2\lambda_\phi v_r^2 & \lambda_{H\phi} v v_r \\ \lambda_{H\phi} v v_r & 2\lambda v^2 \end{bmatrix} \begin{bmatrix} \phi_r \\ h \end{bmatrix}. \quad (5)$$

For the case  $\lambda_{H\phi} v v_r \ll |2\lambda_\phi v_r^2 - 2\lambda v^2|$ , the mixing is suppressed and  $h$  corresponds to  $\phi_r$  with mass  $\simeq 2\lambda v^2$ . The mass of  $\phi_r$  is approximately equal to  $\simeq 2\lambda_\phi v_r^2$ .

Notice that this model in the minimal version shares some features with a model discussed in [5] but in [5] the new scalar degrees of freedom are integrated out. In the present paper, we are more interested in light  $\phi_r$ .

## 2.2 Extended model with two complex scalars

The model in the extended version contains two scalar complex fields  $\phi_1$  and  $\phi_2$  forming a doublet

$$\Phi = \begin{pmatrix} \phi_1 \\ \phi_2 \end{pmatrix} \quad (6)$$

which transforms under the  $U(1)_X$  gauge symmetry as

$$\begin{pmatrix} \phi_1 \\ \phi_2 \end{pmatrix} \rightarrow U \cdot \begin{pmatrix} \phi_1 \\ \phi_2 \end{pmatrix} \quad (7)$$

where

$$U = \begin{pmatrix} \cos \alpha & i \sin \alpha \\ i \sin \alpha & \cos \alpha \end{pmatrix}. \quad (8)$$

Notice that under this transformation

$$\frac{(\phi_1 + \phi_2)}{\sqrt{2}} \rightarrow e^{i\alpha} \frac{(\phi_1 + \phi_2)}{\sqrt{2}}$$

and

$$\frac{(\phi_1 - \phi_2)}{\sqrt{2}} \rightarrow e^{-i\alpha} \frac{(\phi_1 - \phi_2)}{\sqrt{2}}.$$

In other words, in this model we have two complex scalar fields with opposite  $U(1)_X$  charges. However, as we will see, it is more convenient to make the discussion in terms of  $\phi_1$  and  $\phi_2$  forming a doublet representation of  $U(1)_X$ . In addition to the new  $U(1)_X$  gauge symmetry, we also impose a  $Z_2$  symmetry under which all the SM particles are even and the new particles transform as follows:

$$Z_2^{(A)} : \Phi \rightarrow \sigma_3 \Phi \text{ and } V_\mu \rightarrow -V_\mu. \quad (9)$$

Out of the doublet  $\Phi$ , one can make the following bilinear combinations which all are invariant under the gauge symmetry:

$$\Phi^\dagger \Phi = \phi_1^\dagger \phi_1 + \phi_2^\dagger \phi_2 \quad (10)$$

$$\Phi^T \sigma_3 \Phi = \phi_1^2 - \phi_2^2 \quad (11)$$

and

$$\Phi^\dagger \sigma_1 \Phi = \phi_2^\dagger \phi_1 + \phi_1^\dagger \phi_2. \quad (12)$$

Notice that  $\Phi^T \sigma_2 \Phi = 0$ . The combinations in Eq. (10) and (11) are  $Z_2^{(A)}$  even and the one in Eq. (12) is  $Z_2^{(A)}$  odd. Using these combinations, the most general potential involving the scalars of the theory can be written as

$$V(\Phi, H) = -\mu_H^2 H^\dagger H + \lambda_H (H^\dagger H)^2 - \mu^2 \Phi^\dagger \Phi + \lambda (\Phi^\dagger \Phi)^2 \quad (13)$$

$$+ \lambda_{H\Phi} H^\dagger H \Phi^\dagger \Phi + \xi' (\Phi^\dagger \sigma_1 \Phi)^2$$

$$+ [\xi (\Phi^\dagger \Phi) (\Phi^T \sigma_3 \Phi) - \mu'^2 \Phi^T \sigma_3 \Phi + \lambda' (\Phi^T \sigma_3 \Phi)^2 + \lambda'_{H\Phi} H^\dagger H (\Phi^T \sigma_3 \Phi) + \text{h.c.}]$$

where  $H$  denotes the Standard Model Higgs Doublet. Notice that the mass term of form  $\Phi^\dagger \sigma_1 \Phi$  is forbidden by the  $Z_2^{(A)}$  symmetry. The total Lagrangian can be written as

$$\mathcal{L} = \mathcal{L}^{SM} + (D_\mu \Phi)^\dagger (D^\mu \Phi) - V(\Phi, H) - \frac{1}{4} V_{\mu\nu} V^{\mu\nu} \quad (14)$$

where the explicit form of covariant derivative is  $D_\mu = \partial_\mu - ig_V \sigma_1 V_\mu$ . The Vector boson field-strength tensor,  $V_{\mu\nu}$ , is defined as  $V_{\mu\nu} = \partial_\mu V_\nu - \partial_\nu V_\mu$ . Notice that after imposing the  $Z_2^{(A)}$  symmetry as in Eq. (9) (*i.e.*, removing a mass term in the form of  $\Phi^\dagger \sigma_1 \Phi$ ), there will be an additional  $Z_2$  symmetry under which

$$Z_2^{(B)} : \Phi \rightarrow \sigma_3 \Phi \text{ and } V_\mu \rightarrow -V_\mu . \quad (15)$$

For simplicity, we take all the couplings to be real. In our analysis, we make the following conservative assumptions to guarantee the positiveness of the potential at infinity

$$\lambda, \lambda_H, \lambda_{H\phi}, \xi' > 0 \quad \lambda + 2\lambda' > 2|\xi|, \text{ and } \lambda_{H\phi} > 2|\lambda'_{H\phi}| . \quad (16)$$

Moreover, we generally assume  $\lambda_{H\phi}, \lambda'_{H\phi} \lesssim 0.1$  so that the ‘‘SM Higgs’’ still makes sense within this model. Taking  $\mu_H^2$ ,  $\mu^2$  and  $\mu'^2$  (with the convention defined in Eq. 13) positive, both  $H$  and  $\Phi$  will receive vacuum expectation values breaking the electroweak symmetry as well as the  $U(1)_X$  symmetry. Using the freedom of  $SU(2) \times U(1)$  symmetry, we can go to the canonic unitary gauge within which  $H^T = (0 \ (v + h)/\sqrt{2})$ . We can also in general use the global  $U(1)_X$  symmetry to absorb the imaginary component of  $\langle \phi_2 \rangle$  and write  $\Phi^T$  in terms of real components as

$$\Phi^T = \left( \frac{v_r + \phi_r + iv_i + i\phi_i}{\sqrt{2}} \quad \frac{v' + \phi'_r + i\phi'_i}{\sqrt{2}} \right) . \quad (17)$$

A linear combination of these fields will be the massless Goldstone boson which can be absorbed as the longitudinal component of  $V$ . In this gauge, the new vector boson receives a mass equal to  $g_V \sqrt{v_r^2 + v_i^2 + v'^2}$ . The interesting point is that for a significant part of the parameter space, the minimum lies at  $v' = 0$  and  $v_r^2 + v_i^2 \neq 0$ . In this case, the Goldstone boson is

$$G \equiv \frac{-v_i \phi'_r + v_r \phi'_i}{\sqrt{v_i^2 + v_r^2}}$$

which can be absorbed by using the gauge freedom. The combination perpendicular to this combination is a mass eigenstate with nonzero mass:

$$\phi' \equiv \frac{v_r \phi'_r + v_i \phi'_i}{\sqrt{v_i^2 + v_r^2}} .$$

By making a local  $U(1)_X$  transformation,  $G$  can be absorbed and  $\phi_2$  will be of form  $\phi' e^{i\beta}$  where  $\beta = \arctan(v_i/v_r)$ . As long as  $v' = 0$ , the  $Z_2^{(A)}$  symmetry will be preserved, making the lightest particle among  $\phi'$  and the vector boson stable. Thus, the vector can contribute to the dark matter content of the universe if

$$g_V^2(v_r^2 + v_i^2) < m_{\phi'}^2.$$

From now on, we will focus on such a minimum with  $v' = 0$  and we shall assume that  $m_V^2 < m_{\phi'}^2$ . Although, the  $Z_2^{(A)}$  symmetry is maintained, the  $Z_2^{(B)}$  is broken leading to a mixing of the Higgs with  $\phi_r$  and/or  $\phi_i$  as we will discuss below. Throughout this paper, we fix one of the mass eigenvalues  $m_{\delta_3} \simeq m_h$  to 125 GeV as recently announced by the CMS and ATLAS collaborations.

The interaction terms between the gauge boson and scalars are

$$\frac{g_V^2}{2} V_\mu V^\mu [(\phi_i^2 + \phi_r^2 + \phi'^2) + 2(\phi_i v_i + \phi_r v_r)] + \quad (18)$$

$$g_V V^\mu [-\sin \beta (\phi_r \partial_\mu \phi' - \phi' \partial_\mu \phi_r) + \cos \beta (\phi_i \partial_\mu \phi' - \phi' \partial_\mu \phi_i)].$$

Depending on the choice of the parameters of the potential, three phenomenologically distinct regimes can be realized:

- **Phase I**  $v' = 0$ ,  $v_i, v_r \neq 0$ :

In this case both  $\phi_i$  and  $\phi_r$  mix with  $h$ :

$$\begin{pmatrix} \phi_r \\ \phi_i \\ h \end{pmatrix} = \begin{pmatrix} a_{11} & a_{12} & a_{13} \\ a_{21} & a_{22} & a_{23} \\ a_{31} & a_{32} & a_{33} \end{pmatrix} \begin{pmatrix} \delta_1 \\ \delta_2 \\ \delta_3 \end{pmatrix} \quad (19)$$

where  $\delta_i$  are the mass eigenstates. As a result, both  $\phi_r$  and  $\phi_i$  become unstable. Notice that in this case, CP is broken spontaneously. The values of  $a_{ij}$  in terms of the parameters of the potential are given in the appendix. We assume that the mixings between the SM Higgs and the new scalars are small:  $a_{31}, a_{32}, a_{13}, a_{23} \ll 1$ .

- **Phase II**  $v' = v_r = 0$  and  $v_i \neq 0$ ;

In this case only  $\phi_i$  mixes with  $h$ . That is in the matrix shown in Eq. (19),  $a_{12} = a_{13} = a_{21} = a_{31} = 0$ . Taking  $\phi_i$  to be CP-even, CP will be preserved. For  $v_r = 0$ , the Lagrangian is invariant under  $\phi_r \rightarrow -\phi_r$ .

Thus, in addition to  $Z_2^{(A)}$ , there is another  $Z_2$  which preserves  $\delta_1 = \phi_r$  against decay. As a result, there will be two candidates for dark matter:  $\delta_1 = \phi_r$  and  $V$ . In this phase,  $\phi'_r$  is the Goldstone boson which can be absorbed (*i.e.*,  $\beta = \pi/2$ ) so we have  $\phi' = \phi'_i$ .

- **Phase III**  $v' = v_i = 0$  and  $v_r \neq 0$ ;

In this case only  $\phi_r$  mixes with  $h$ . That is in the matrix shown in Eq. (19),  $a_{12} = a_{23} = a_{21} = a_{32} = 0$ . CP will be preserved. In this phase,  $\phi'_i$  is the Goldstone boson which can be absorbed (*i.e.*,  $\beta = 0$ ) so we have  $\phi' = \phi'_r$ . It is straightforward to verify that the phases II and III are equivalent provided that we substitute

$$(\mu^2, \lambda_H, \lambda, \xi', \lambda'_{H\phi}, \mu'^2, \lambda'_{H\phi}, \xi) \rightarrow (\mu^2, \lambda_H, \lambda, \xi', \lambda', \lambda_{H\phi}, -\mu'^2, -\lambda'_{H\phi}, -\xi)$$

and

$$\phi'_i \leftrightarrow \phi'_r .$$

In all these phases, in the “unitary” gauge where the Goldstone boson is absorbed,  $V_\mu$  is a space-like vector with three polarizations satisfying  $\partial^\mu V_\mu = 0$ .

In the appendix, we have formulated the conditions on the parameters of the model under which each of these cases are realized.

### 3 Annihilation of dark matter pair

In the previous section, we introduced the models with one and two scalar fields. In this section, we discuss how the new stable particles can account for the dark matter content of the universe for each case one by one.

#### 3.1 Annihilation modes in the minimal model

The scalar field,  $\phi_r$ , can be either lighter or heavier than  $V_\mu$ . If it is heavier, the main annihilation mode for the  $V$  pair will be through  $s$ -channel scalar exchange. Taking the  $\phi_r - h$  mixing small, we find

$$\langle \sigma(V + V \rightarrow \text{final}) v_{rel} \rangle = \frac{64}{3} g_V^4 \left[ \frac{\lambda_{H\phi} v v'}{(m_h^2 - 4m_V^2)(m_{\phi_r}^2 - 4m_V^2)} \right]^2 F \quad (20)$$



with

$$F \equiv \lim_{m_{h^*} \rightarrow 2m_V} \left( \frac{\Gamma(h^* \rightarrow final)}{m_{h^*}} \right). \quad (21)$$

Here  $\Gamma(h^* \rightarrow final)$  denotes the rate for the decay mode,  $h^* \rightarrow final$ , for a hypothetical SM-like Higgs,  $h^*$ , whose mass is  $m_{h^*} = 2m_V$ . For  $m_b < m_{DM} < m_W, m_{\phi_r}$ , the main annihilation mode will be to a  $b\bar{b}$  pair [13] which is constrained by bounds on the antiproton flux from PAMELA [14, 15]. By adding a new Higgs doublet exclusively coupled to leptons, the bound can be circumvented but we shall not discuss this possibility. Let us now discuss the case that  $m_W < m_V < m_{\phi_r}$ . In this case, the main annihilation mode is to the  $W^+W^-$  pair. If  $m_Z < m_V$ , the dark matter can also annihilate to a  $Z$  pair with

$$\frac{\sigma(V + V \rightarrow Z + Z)}{\sigma(V + V \rightarrow W^+ + W^-)} = \frac{\text{Br}(H^* \rightarrow Z + Z)}{\text{Br}(H^* \rightarrow W^+ + W^-)} \Big|_{m_{H^*}=2m_V} < 1.$$

The subsequent decay of  $W^+$  and  $W^-$  can produce detectable secondary particles. In particular, the bounds from the antiproton and gamma ray fluxes are strong. The present bound from PAMELA on antiproton flux [14] as well as the gamma ray bound from Fermi-LAT [16] lie above 1 pb [10]. However, the forthcoming AMS02 experiment [17] may be able to probe this scenario. Notice that in this range of parameters, the model shares some features with the model discussed in [10] with the difference that here  $g_V$  is a free parameter independent of the gauge interactions of the standard model so a wider range of  $m_{DM}$  is possible.

Let us now discuss the case that  $m_{\phi_r} < m_{DM}$ . In this case, new annihilation mode for a pair of  $V$  become possible:

$$\langle \sigma(V + V \rightarrow \phi_r + \phi_r) v_{rel} \rangle = \frac{g_V^4}{24\pi m_V^2} g(m_{\phi_r}^2/m_V^2) \quad (22)$$

where

$$g(x) = \sqrt{1-x} \left( \left(1 + \frac{4}{x-2}\right)^2 + \frac{16}{3} \frac{(1-x)^2}{(x-2)^2} + \frac{8}{3} \left(\frac{1-x}{x-2}\right) \left(1 + \frac{4}{x-2}\right) \right).$$

The produced  $\phi_r$  are unstable and will eventually decay to the SM particles. If  $2m_b < m_{\phi_r} < 2m_W$ , once the  $V$  pair annihilate to the  $\phi_r$  pair, the annihilation products will eventually decay to a  $b\bar{b}$  pair leading to an excess in

antiproton flux which is restricted by PAMELA. In the following two ranges, the antiproton bounds can be avoided: 1)  $2m_W < m_{\phi_r} < m_V$ ; 2)  $m_V \sim \text{few GeV}$  and  $m_{\phi_r} < 2m_p$ . An example of the first range is

Point I :  $m_V = 250 \text{ GeV}$ ,  $m_{\phi_r} = 200 \text{ GeV}$ ,  $v' = 1023 \text{ GeV}$ ,  $\lambda_\phi = 0.13$ ,  $g_V = 0.24$

and examples of the second range are

Point II :  $m_V = 8 \text{ GeV}$ ,  $m_{\phi_r} = 1.5 \text{ GeV}$ ,  $v' = 187 \text{ GeV}$ ,  $\lambda_\phi = 0.005$ ,  $g_V = 0.042$

or

Point III :  $m_V = 10 \text{ GeV}$ ,  $m_{\phi_r} = 1.5 \text{ GeV}$ ,  $v' = 210 \text{ GeV}$ ,  $\lambda_\phi = 0.005$ ,  $g_V = 0.047$

where we have taken the mixing given by  $\lambda_{H\phi}$  to be small.

### 3.2 Annihilation modes in the extended model

As we discussed in section 2.2, in order for  $V_\mu$  to be stable, it should be lighter than  $\phi'$ . However, the other scalars  $\delta_i$  (mass eigenstates made of  $\phi_i$ ,  $\phi_r$  and  $h$ ) can be either lighter or heavier. In case that  $\delta_i$  are all heavier than  $V_\mu$ , the model can be considered as a Higgs portal model within which

$$\langle \sigma(V + V \rightarrow \text{final}) v_{rel} \rangle = \frac{64}{3} g_V^4 \left[ \sum_{j=1}^3 \frac{a_{3j}(a_{1j}v_r + a_{2j}v_i)}{m_{\delta_j}^2 - 4m_V^2} \right]^2 F \quad (23)$$

where  $F$  is defined in Eq. (21). The rest of the discussion is as the case with a single scalar.

Let us now discuss the case that  $m_{\delta_1} < m_{DM}$ . Like the case of a single scalar, in this case too, new annihilation mode(s) for a pair of  $V$  become possible.

The interaction terms within the term  $(D_\mu \Phi \cdot D^\mu \Phi)$  lead to the Feynman diagrams shown in Fig. 1. We denote the amplitude of diagram in Fig. (1-a) with  $M_1$  and the sum of the amplitudes of diagrams in Figs. (1-b) and (1-c) with  $M_2$ . We denote the sum of the amplitudes of diagrams in Figs (1-d) and (1-e) with  $M_3$ . There will be also a contribution from trilinear couplings  $A_j \delta_1^2 \delta_j$  which comes from the quartic terms of the potential in Eq. (13) as shown in Fig. 2. We denote the amplitude of diagram in Fig. 2 with  $M_4$ .

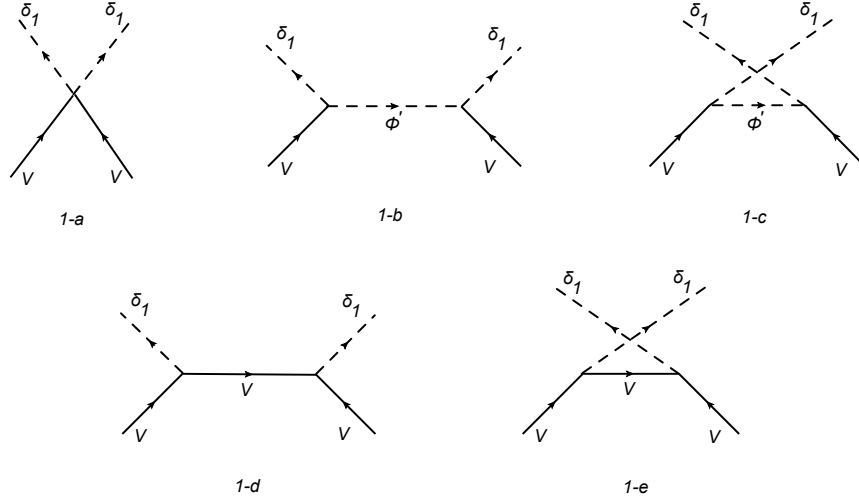


Figure 1: Feynman diagrams for the annihilation of the DM pair to a  $\delta_1$  pair

The total annihilation cross section is

$$\begin{aligned}
 \langle \sigma_{tot}(V + V \rightarrow \delta_1 \delta_1) v_{rel} \rangle &= \frac{\sqrt{1 - \frac{m_{\delta_1}^2}{m_V^2}}}{64\pi m_V^2} \frac{1}{9} \sum_{spins} |M_1 + M_2 + M_3 + M_4|^2 \\
 &= \sum_{n=1(n \geq m)}^4 \sum_{m=1}^4 \sigma_{mn}
 \end{aligned} \tag{24}$$

where  $\sigma_{mn}$  come from the interference of  $M_m$  and  $M_n$ :

$$\begin{aligned}
 \sigma_{11} &= \frac{g_V^4 \sqrt{1 - \frac{m_{\delta_1}^2}{m_V^2}}}{48\pi m_V^2} (|a_{11}|^2 + |a_{21}|^2)^2 \\
 \sigma_{12} &= \frac{2g_V^4 \sqrt{1 - \frac{m_{\delta_1}^2}{m_V^2}}}{9\pi m_V^2} \frac{(m_V^2 - m_{\delta_1}^2) S^2 (|a_{11}|^2 + |a_{21}|^2)}{(m_{\phi'}^2 + m_V^2 - m_{\delta_1}^2)} \\
 \sigma_{22} &= \frac{16g_V^4 \sqrt{1 - \frac{m_{\delta_1}^2}{m_V^2}}}{9\pi m_V^2} \frac{S^4 (m_V^2 - m_{\delta_1}^2)^2}{(m_{\phi'}^2 + m_V^2 - m_{\delta_1}^2)^2} \\
 \sigma_{13} &= \frac{g_V^6 \sqrt{1 - \frac{m_{\delta_1}^2}{m_V^2}}}{18\pi m_V^4} (|a_{11}|^2 + |a_{21}|^2) (a_{11} v_r + a_{21} v_i)^2 \frac{m_{\delta_1}^2 - 4m_V^2}{2m_V^2 - m_{\delta_1}^2}
 \end{aligned} \tag{25}$$

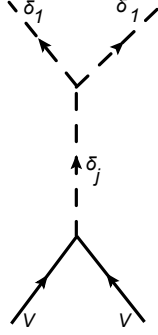


Figure 2: The  $s$ -channel contribution to the DM pair annihilation

$$\begin{aligned}
\sigma_{23} &= \frac{8g_V^6(1 - \frac{m_{\delta_1}^2}{m_V^2})^{3/2}}{9\pi m_V^2} \frac{S^2(a_{11}v_r + a_{21}v_i)^2}{m_{\phi'}^2 + m_V^2 - m_{\delta_1}^2} \\
\sigma_{33} &= \frac{g_V^8 \sqrt{1 - \frac{m_{\delta_1}^2}{m_V^2}}}{9\pi m_V^2} \frac{(a_{11}v_r + a_{21}v_i)^4}{(2m_V^2 - m_{\delta_1}^2)^2} [6 - 4\frac{m_{\delta_1}^2}{m_V^2} + \frac{m_{\delta_1}^4}{m_V^4}] \\
\sigma_{14} &= \frac{g_V^4 \sqrt{1 - \frac{m_{\delta_1}^2}{m_V^2}}}{12\pi m_V^2} [\sum_{j=1}^3 \frac{A_j(a_{1j}v_r + a_{2j}v_i)}{m_{\delta_j}^2 - 4m_V^2}] (|a_{11}|^2 + |a_{21}|^2) \\
\sigma_{24} &= \frac{4S^2 g_V^4 \sqrt{1 - \frac{m_{\delta_1}^2}{m_V^2}}}{9\pi m_V^2} [\sum_{j=1}^3 \frac{A_j(a_{1j}v_r + a_{2j}v_i)}{m_{\delta_j}^2 - 4m_V^2}] \frac{m_V^2 - m_{\delta_1}^2}{m_{\phi'}^2 + m_V^2 - m_{\delta_1}^2} \\
\sigma_{34} &= \frac{g_V^6 \sqrt{1 - \frac{m_{\delta_1}^2}{m_V^2}}}{9\pi m_V^4} [\sum_{j=1}^3 \frac{A_j(a_{1j}v_r + a_{2j}v_i)}{m_{\delta_j}^2 - 4m_V^2}] \frac{(a_{11}v_r + a_{21}v_i)^2 (4m_V^2 - m_{\delta_1}^2)}{2m_V^2 - m_{\delta_1}^2} \\
\sigma_{44} &= \frac{g_V^4 \sqrt{1 - \frac{m_{\delta_1}^2}{m_V^2}}}{12\pi m_V^2} [\sum_{j=1}^3 \frac{A_j(a_{1j}v_r + a_{2j}v_i)}{m_{\delta_j}^2 - 4m_V^2}]^2,
\end{aligned} \tag{26}$$

in which  $S^2 = |a_{11}|^2 \sin^2 \beta + 2\Re[a_{11}a_{21}^*] \sin \beta \cos \beta + |a_{21}|^2 \cos^2 \beta$ . In case  $m_{\delta_i} + m_{\delta_k} < 2m_V$ , the pair of  $V$  can annihilate to  $\delta_i + \delta_k$ . The annihilation

rate will be given by the same formula with replacement  $a_{11}^2 \rightarrow a_{1i}a_{1k}$ ,  $a_{21}^2 \rightarrow a_{2i}a_{2k}$  and  $A_j$  with the couplings of  $\delta_j\delta_k\delta_i$ .

Within phase I, all  $\delta_i$  are unstable and the discussion is similar to the case of minimal model with only one scalar. However phases II and III will have a totally different phenomenology. As we saw earlier, phase III is equivalent to phase II. In the following, we focus specifically on the novelties of phase II which applies to phase III, too. Similarly to phase I, in case that  $\delta_1(=\phi_r)$  is lighter than  $V$ , a pair of  $V$  can annihilate to  $\delta_1$ . This annihilation mode will dominate over the annihilation to the SM particles. The cross section is given by the formulas in Eq. (25) setting  $v_r = \cos\beta = 0$ . Remember that in this phase  $\delta_1(=\phi_r)$  is also stable and a component of dark matter. The  $\delta_1$  pair will annihilate via Higgs portal with a cross section

$$\langle\sigma(\delta_1 + \delta_1 \rightarrow h^* \rightarrow \text{final})v_{rel}\rangle = \frac{2(\lambda_{H\phi} + 2\lambda'_{H\phi})^2 v_H^2}{(4m_{\delta_1}^2 - m_h^2)^2} F' , \quad (27)$$

where  $F' \equiv \lim_{m_{h^*} \rightarrow 2m_{\delta_1}} \left( \frac{\Gamma(h^* \rightarrow \text{final})}{m_{h^*}} \right)$ . An interesting scenario is the case that  $m_{\delta_1} < m_V < m_{\delta_2}, m_{\phi'}$ ,

$$\sigma(V + V \rightarrow \text{anything}) = 1 \text{ pb}$$

and

$$\sigma(\delta_1 + \delta_1 \rightarrow h^* \rightarrow \text{final}) \gg 1 \text{ pb}.$$

Taking  $\lambda_{H\phi} + 2\lambda'_{H\phi} \sim 0.1$  this scenario can be realized for *e.g.*, the point in table 1. In this case, at the time of the decoupling of  $V$ ,  $\delta_1$  will be in thermal equilibrium with the SM particles. Since the annihilation cross section of  $\delta_1$  is large, its density will be suppressed so the dark matter will be mainly composed of the  $V$  particles. The interesting point is that here even if  $2m_b < m_{\delta_1} < 2m_W$ , the antiproton bound does not rule out the model as the annihilation of the  $V$  pair will lead to stable  $\delta_1$  and the density of  $\delta_1$  in the present time will be too low for its annihilation to lead to a sizeable secondary flux. For the case that the annihilation cross section of  $\delta_1$  is comparable to that of  $V$ , both particles will have comparable densities at the present time. In case that  $m_{\delta_1} > m_V$ , the  $\delta_1\delta_1$  pair can annihilate to the  $V$  pair. Discussing all these options is beyond the scope of the present paper and will be done elsewhere.

Table 1: Parameters of model. An example for phase II with  $m_{\delta_1} < m_V$ .

$\xi$	$\lambda'$	$\xi'$	$\lambda$	$\mu$ (GeV)	$\mu'$ (GeV)	$g_V^2$	$\lambda_{H\phi}$	$\lambda'_{H\phi}$
0.5	0.11	0.4	0.93	409	146	0.017	0.1	0.1

Table 2: The mass spectrum of the point in table 1. All the masses are in GeV.

$v_i$	$v_r$	$m_{\delta_1}$	$m_{\delta_2}$	$m_{\phi'}$	$m_V$
893	0	100	500	1000	116

## 4 Direct DM detection and signature at the colliders

In this section, we will discuss the discovery potential of the model via direct dark matter searches and the production at the colliders. Since the only portal between the new sector and the SM particles is through the SM Higgs, both processes are determined by the couplings  $\lambda_{H\phi}$  and  $\lambda'_{H\phi}$ . Let us first discuss the bound from direct searches. The interaction of  $V$  and the nucleon is through Higgs portal and is therefore spin-independent. Throughout this discussion, we assume that the vector boson is the main component of the dark matter. Translating the results of [7] for vector field in case of single scalar fermion, we find

$$\sigma_N \equiv \sigma_{SI}(V + N \rightarrow V + N) = \frac{g_V^4 M_r^2 m_N^2}{\pi m_V^2 v_H^2} \left[ \frac{\lambda_{H\phi} v v_r^2}{m_h^2 m_{\phi_r}^2} \right]^2 f^2 \quad (28)$$

where  $M_r = (m_V^{-1} + m_N^{-1})^{-1}$  is the DM-nucleon reduced mass and  $0.14 < f < 0.66$ . Notice that we have taken the mixing between  $h$  and  $\phi_r$  small. For two scalar model, we find

$$\sigma_N \equiv \sigma_{SI}(V + N \rightarrow V + N) = \frac{g_V^4 M_r^2 m_N^2}{\pi m_V^2 v_H^2} \left[ \left( \sum_{j=1}^3 \frac{a_{3j}(a_{1j} v_r + a_{2j} v_i)}{m_{\delta_j}^2} \right) \right]^2 f^2 \quad (29)$$

For the phase II, Eq. (29) simplifies as

$$\sigma_N = (\lambda_{H\phi} - 2\lambda'_{H\phi})^2 \frac{m_V^2 M_r^2 m_N^2}{\pi m_{\delta_2}^4 m_h^4} f^2. \quad (30)$$

For phase III, the formula for scattering cross section in Eq. (30) also applies after replacing  $m_{\delta_2} \rightarrow m_{\delta_1}$  and  $\lambda'_{H\phi} \rightarrow -\lambda'_{H\phi}$ .

In the following, we focus on the model with two scalars. The discussion of the model with a single scalar is very similar to that of phase I if we replace  $\delta_1$  with  $\phi_r$  and  $a_{31}$  with  $\lambda_{H\phi} v v_r / (2\lambda_\phi v_r^2 - 2\lambda v^2)$ . Of course in the minimal version there is no  $\lambda'_{H\phi}$  coupling and no  $\phi'$  or  $\delta_2$  field.

Let us now discuss the signature at the high energy colliders. At the colliders,  $\delta_i$  can be in principle produced with a cross section that is suppressed by  $|a_{3i}|^2$  relative to the SM Higgs of mass  $m_{\delta_i}$ . For  $m_{\delta_i} > 2m_W$  and relatively large  $|a_{3i}|^2$ , this may provide an observable signal. In this case, we expect a SM-like Higgs decaying to  $W^+W^-$  but with a production rate suppressed by a factor of  $a_{3i}^2 \sim 0.1$ . The LHC should be able to put a bound on  $|a_{3i}|^2$  for this range. Now, let us discuss another potential discovery channel. If  $\lambda_{H\phi}$  and  $\lambda'_{H\phi}$  are relatively large and some of the new neutral scalar particles are lighter than  $m_h/2$ , the Higgs should have sizeable invisible decay modes. For simplicity let us take  $\lambda'_{H\phi} \ll \lambda_{H\phi}$ . The invisible decay rate of the Higgs to first order of approximation will be  $\sim \lambda_{H\phi}^2 v_H^2 / (64\pi m_h)$ . In fact, the bound from the LHC on the Higgs invisible decay rate [11] already rules out  $\lambda_{H\phi} > 0.01$  for  $m_{\delta_{1,2}} < m_h/2 \simeq 63$  GeV.

Let us now discuss each of the scenarios of interest one by one.

- **Phase I with  $2m_W < m_{\delta_1} < m_V$**

Within this scenario  $\delta_i$  are too heavy to lead to invisible decay width for the Higgs. However, if  $\lambda_{H\phi}, \lambda'_{H\phi} > 0.1$  we might observe two SM-like Higgses decaying to  $W^+W^-$  pair at masses of  $m_{\delta_1}$  and  $m_{\delta_2}$ . For such values of  $\lambda_{H\phi}$  and  $\lambda'_{H\phi}$ , the scattering cross section off nuclei is of order of  $10^{-46} \lambda_{H\phi}^2 (f/0.3)^2 \text{ cm}^2$  which is well below the bound from XENON100 [12].

- **Phase I with  $m_{\delta_1} < m_V$  and  $m_{\delta_1} < 2m_p$**

At this range, there are light bosons so there is a possibility of  $\delta_3 (\simeq h) \rightarrow \delta_1 + \delta_1, \delta_2 + \delta_2, \delta_1 + \delta_2, \phi' + \phi'$ . The rate depends on the values of  $\lambda_{H\phi}, \lambda'_{H\phi}$ . The scattering cross section is expected to be of order of

$10^{-42} \lambda_{H\phi}^2 (f/0.3)^2 \text{ cm}^2$  which is well below the current bounds [12] even for large  $\lambda_{H\phi}$  and  $\lambda'_{H\phi}$ .

- **Phase II with  $m_{\delta_1} < m_V$**

Here,  $\delta_1$  is stable. If the scalars are heavier than  $m_h/2$ , the Higgs will not have an invisible decay mode. The interesting point here is that we require  $\lambda_{H\phi}$  and  $\lambda'_{H\phi}$  to be relatively large from constraint on the  $\delta_1$  density. As result, the model can be tested by searching for an excess in  $W^+W^-$  pair with an invariant mass equal to  $m_{\delta_2}$ . From the bound on the invisible decay width of the Higgs [11], we already know  $m_{\delta_1} > 63 \text{ GeV}$ . For such values of  $\lambda_{H\phi}$  and  $\lambda'_{H\phi}$  and for  $f = 0.3$ , the scattering cross section is about  $10^{-45} - 10^{-46} \text{ cm}^2$  which is below the current XENON100 bound [12].

## 5 Lower bound on coupling of new scalars with Higgs

Let us first focus on the model with two scalars. At the limit  $\lambda_{H\phi}, \lambda'_{H\phi} \rightarrow 0$ , the dark sector including  $\Phi$  and  $V$  will decouple from the SM sector. As we saw within phase II,  $\lambda_{H\phi}$  and  $\lambda'_{H\phi}$  should be relatively large so that the lighter dark matter component ( $\delta_1$  for the point in table 1 and 2) has sufficiently large annihilation cross section. However, there is no such a constraint for the minimal model or the phase I of the extended model. A lower bound on  $\lambda_{H\phi}, \lambda'_{H\phi}$  comes from the assumption that DM is produced thermally in the early universe through interaction with the Higgs. The production rate at high temperature is expected to be given by  $O[(\lambda_{H\phi}^2 + \lambda_{H\phi}'^2)T/(4\pi)]$  which should be compared with the Hubble expansion rate at the time. Setting the ratio of the production rate to the Hubble constant at  $T = m_{\delta_i}$  larger than one, we find

$$\text{Max}[\lambda_{H\phi}, \lambda'_{H\phi}] \gtrsim 10^{-8} \left( \frac{m_{\delta}}{100 \text{ GeV}} \right)^{1/2}.$$

The lower bound from the decay of the unstable new particles before big bang nucleosynthesis epoch is much weaker.

The discussion for the minimal model with a single scalar is similar provided that we replace  $\delta_i$  with  $\phi_r$  and drop  $\lambda'_{H\phi}$ .



## 6 Conclusions

We have introduced a simple model based on a new  $U(1)_X$  gauge symmetry within which the vector boson plays the role of the dark matter. The model also contains new complex scalar(s). In the minimal version, there is only one scalar and in the extended version there are two scalars:  $\phi_1$  and  $\phi_2$ . One of these scalar fields ( $\phi_1$  for the extended version) develops a Vacuum Expectation Value (VEV) which breaks the  $U(1)_X$  symmetry. The dark matter is protected from decay by a  $Z_2$  symmetry. The coupling between the new sector and the SM sector is through the scalar couplings with the Higgs:  $\lambda_{H\phi}$  and  $\lambda'_{H\phi}$ .

In the extended version, depending on the range of parameters, different phases with different phenomenology appear. In one phase, both  $\text{Im}(\phi_1)$  and  $\text{Re}(\phi_1)$  receive a VEV leading to spontaneous CP-violation. In another phase, only one of  $\text{Im}(\phi_1)$  and  $\text{Re}(\phi_1)$  receives a VEV. In the latter case, CP is conserved. Moreover, a new  $Z_2$  appears that leads to a new stable dark matter candidate. From the abundance of the lighter dark matter candidate we obtain a lower bound on the couplings of new scalar with the Higgs (*i.e.*, on  $\lambda_{H\phi}$  and  $\lambda'_{H\phi}$ ) of order of 0.1.

We have briefly discussed the possibility of discovery at the colliders and direct dark matter searches. At colliders, the new particles can show up as invisible decay mode of the SM Higgs or as extra SM-like Higgs(es) with a suppressed production rate at various masses. The scattering cross section off nuclei is expected to be at least one order of magnitude below the present bound. Within the phase II of the extended model, the upper bound on the invisible decay width of the Higgs [11] implies that the masses of the new particles are larger than  $m_h/2 = 63$  GeV.

## 7 Appendix A

As discussed earlier when  $v_i^2 + v_r^2 \neq 0$ , the Higgs field mixes with  $\phi_i$  and/or  $\phi_r$  defined in Eq. (17). The components of the mass matrix in the  $(\phi_r \ \phi_i \ h)$  basis are

$$\begin{aligned} M_{11}^2 &= -\mu^2 - 2\mu'^2 + (\lambda - 6\lambda')v_i^2 + 3(\lambda + 2\lambda' + 2\xi)v_r^2 + (\lambda_{H\phi} + 2\lambda'_{H\phi})\frac{v_H^2}{2} \\ M_{22}^2 &= -\mu^2 + 2\mu'^2 + 3(\lambda + 2\lambda' - 2\xi)v_i^2 + (\lambda - 6\lambda')v_r^2 + (\lambda_{H\phi} - 2\lambda'_{H\phi})\frac{v_H^2}{2} \end{aligned}$$

$$\begin{aligned}
M_{33}^2 &= -\mu_H^2 + 3\lambda_H v_H^2 + \frac{\lambda_{H\phi}}{2}(v_i^2 + v_r^2) + \lambda'_{H\phi}(v_r^2 - v_i^2) \\
M_{12}^2 &= 2(\lambda - 6\lambda')v_i v_r \\
M_{13}^2 &= (\lambda_{H\phi} + 2\lambda'_{H\phi})v_H v_r \\
M_{23}^2 &= (\lambda_{H\phi} - 2\lambda'_{H\phi})v_H v_i .
\end{aligned}$$

The components of mass matrix of  $(\phi'_r, \phi'_i)$  are

$$\begin{aligned}
M_{rr}^2 &= -\mu^2 + 2\mu'^2 + v_i^2(\lambda + 2\lambda' - 2\xi) + v_r^2(\lambda - 2\lambda' + 2\xi') + \frac{v_H^2}{2}(\lambda_{H\phi} - 2\lambda'_{H\phi}) \\
M_{ii}^2 &= -\mu^2 - 2\mu'^2 + v_i^2(\lambda - 2\lambda' + 2\xi') + v_r^2(\lambda + 2\lambda' + 2\xi) + \frac{v_H^2}{2}(\lambda_{H\phi} + 2\lambda'_{H\phi}) \\
M_{ir}^2 &= 2(\xi' + 2\lambda')v_i v_r .
\end{aligned} \tag{31}$$

In the limit  $\lambda_{H\phi}, \lambda'_{H\phi} \ll 1$ , we can write  $m_h^2 \simeq M_{33}^2$  and the mixing matrix in Eq. (19) approximately as

$$\begin{bmatrix} \cos \theta & \sin \theta & a_{13} \\ -\sin \theta & \cos \theta & a_{23} \\ a_{31} & a_{32} & 1 \end{bmatrix} \tag{32}$$

where

$$\tan 2\theta = \frac{2M_{12}^2}{M_{22}^2 - M_{11}^2}$$

and

$$\begin{aligned}
a_{13} &= \frac{(\lambda_{H\phi} S_1 + 2\lambda'_{H\phi} S_2)v_H}{(m_h^2 - m_{\delta_1}^2)(m_h^2 - m_{\delta_2}^2)} \\
a_{23} &= \frac{(\lambda_{H\phi} T_1 + 2\lambda'_{H\phi} T_2)v_H}{(m_h^2 - m_{\delta_1}^2)(m_h^2 - m_{\delta_2}^2)}
\end{aligned} \tag{33}$$

in which

$$\begin{aligned}
S_1 &= -M_{33}^2(\cos \theta v_r + \sin \theta v_i) + M_{22}^2 v_r \cos \theta + M_{11}^2 v_i \sin \theta - M_{12}^2(v_r \sin \theta + v_i \cos \theta) \\
S_2 &= -M_{33}^2(\cos \theta v_r - \sin \theta v_i) + M_{22}^2 v_r \cos \theta - M_{11}^2 v_i \sin \theta - M_{12}^2(v_r \sin \theta - v_i \cos \theta) \\
T_1 &= M_{33}^2(\sin \theta v_r - \cos \theta v_i) - M_{22}^2 v_r \sin \theta + M_{11}^2 v_i \cos \theta - M_{12}^2(v_r \cos \theta - v_i \sin \theta) \\
T_2 &= M_{33}^2(\sin \theta v_r + \cos \theta v_i) - M_{22}^2 v_r \sin \theta - M_{11}^2 v_i \cos \theta - M_{12}^2(v_r \cos \theta + v_i \sin \theta).
\end{aligned} \tag{34}$$

Finally,

$$\begin{aligned} a_{31} &= a_{23} \sin \theta - a_{13} \cos \theta \\ a_{32} &= -a_{23} \cos \theta - a_{13} \sin \theta . \end{aligned} \quad (35)$$

In phase I, we have

$$\begin{aligned} v_r^2 &= \frac{\mu^2(4\lambda' - \xi) + 2\mu'^2(\lambda - 2\lambda' - 2\xi)}{2(4\lambda\lambda' - 8\lambda'^2 - \xi^2)}, \\ v_i^2 &= \frac{\mu^2(4\lambda' + \xi) + 2\mu'^2(-\lambda + 2\lambda' - \xi)}{2(4\lambda\lambda' - 8\lambda'^2 - \xi^2)}, \end{aligned}$$

and

$$m_{\phi'}^2 = \frac{4(2\lambda' + \xi')(2\lambda'\mu^2 - \xi\mu'^2)}{4\lambda\lambda' - 8\lambda'^2 - \xi^2}. \quad (36)$$

Of course, in order for the phase I to be realized, the parameters have to be in a range where  $v_i^2, v_r^2, m_{\phi'}^2 > 0$ . The stability of  $V$  (*i.e.*,  $m_V^2 = g_V^2(v_i^2 + v_r^2) < m_{\phi'}^2$ ) then implies

$$g_V < \sqrt{2(2\lambda' + \xi')}. \quad (37)$$

In order for the phase II with  $v_r = v' = 0$  to be realized, we should have

$$\mu^2 - 2\mu'^2 > 0, \quad \mu^2(\xi + \xi' - 2\lambda') + 2\mu'^2(\xi + \xi' - \lambda) > 0$$

and

$$\mu^2(\xi - 4\lambda') + 2\mu'^2(\xi + 2\lambda' - \lambda) > 0.$$

In phase II, we find  $m_{\delta_1}^2 = M_{11}^2$ ,  $m_{\delta_2}^2 \simeq M_{22}^2$ ,  $m_h^2 \simeq m_{\delta_3}^2 \simeq M_{33}^2$  and  $a_{12} = a_{21} = a_{13} = a_{31} = 0$  and

$$a_{23} = -a_{32} = \frac{(\lambda_{H\phi} - 2\lambda'_{H\phi})v_i v_H}{m_{\delta_2}^2 - m_h^2}.$$

Within the phase II

$$v_i^2 = \frac{\mu^2 - 2\mu'^2}{\lambda + 2\lambda' - 2\xi}$$

and

$$m_{\phi'}^2 = \frac{\mu^2(2\xi + 2\xi' - 4\lambda') + \mu'^2(4\xi - 4\lambda - 4\xi')}{\lambda + 2\lambda' - 2\xi}. \quad (38)$$

The stability of  $V$  (*i.e.*,  $m_V^2 = g_V^2(v_i^2 + v_r^2) < m_{\phi'}^2$ ) then implies

$$g_V < \left( \frac{2[\mu^2(\xi + \xi' - 2\lambda') + \mu'^2(2\xi - 2\xi' - 2\lambda)]}{\mu^2 - 2\mu'^2} \right)^{1/2}. \quad (39)$$

In phase III where  $v_r \neq 0$  and  $v_i = 0$ , we find  $m_{\delta_2}^2 = M_{22}^2$ ,  $m_{\delta_1}^2 \simeq M_{11}^2$ ,  $m_h^2 \simeq m_{\delta_3}^2 \simeq M_{33}^2$  and  $a_{12} = a_{21} = a_{23} = a_{32} = 0$  and

$$a_{13} = -a_{31} = \frac{(\lambda_{H\phi} + 2\lambda'_{H\phi}2)v_r v_H}{m_{\delta_1}^2 - m_h^2}.$$

Within this phase

$$v_r^2 = \frac{\mu^2 + 2\mu'^2}{\lambda + 2(\lambda' + \xi)} \quad (40)$$

and

$$m_{\phi'}^2 = \frac{2\mu^2(\xi' - \xi - 2\lambda') + 4\mu'^2(\xi + \xi' + \lambda)}{\lambda + 2(\lambda' + \xi)}. \quad (41)$$

The stability of  $V$  (*i.e.*,  $m_V^2 = g_V^2(v_i^2 + v_r^2) < m_{\phi'}^2$ ) then implies

$$g_V < \left( \frac{2[\mu^2(\xi' - \xi - 2\lambda') + 2\mu'^2(\lambda + \xi + \xi')]}{\mu^2 + 2\mu'^2} \right)^{1/2}. \quad (42)$$

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